# THE DYNAMICS OF CHARGES INDUCED BY A CHARGED PARTICLE TRAVERSING A DIELECTRIC SLAB

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**Abstract**

We studied the dynamics of surfacea and wake charges induced by a charged particle traversing a dielectric slab. It is shown that after the crossing of the slab first boundary, the induced on the slab surface charge (image charge) is transformed into the wake charge, which overflows to the second boundary when the particle crosses it. It is also shown, that the polarization of the slab is of an oscillatory nature, and the net induced charge in a slab remains zero at all stages of the motion.

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## 1 Introduction

As it passes through a medium, a fast charged particle excites oscillations of the charge density behind itself [1-3]. These wakefields and the particle energy losses associated with their excitation have been studied widely for a variety of media [3-8]. Wakefields have recently reattracted interest because of the development of new methods for accelerating particles [9-11].

In most studies of wakefields it has been assumed that the medium is unbounded. The wakefields are excited as the particle enters the medium, or they disappear when the particle leaves the medium, because of various transient polarization processes which occur near the interface. Among these processes, the excitation of surface oscillations and the associated additional energy loss have been studied previously [12-20]. In connection with the development of new particle acceleration methods, numerical calculations have determined the distance from the sharp plasma boundary at which the amplitude of the wakefield excited by an ultrarelativistic particle reaches the same level as in an unbounded medium [21].

Fairly recently, in connection with problems of emission electronics and optoelectronics, the image charge and the dynamical image potential created by a moving particle has also been investigated. To describe the process of formation of the image charge, various approaches (quantum mechanical, the hydrodynamic, etc.) and various models of the medium have been employed [22-25].

In the present paper we analyze the dynamics of reversal of the sign of the charges induced at the slab boundary (repolarization of the slab) as the particle crosses the interface. The process is found to be of a nonmonotonic, oscillatory nature. The case of normal incidence of a particle through the slab is considered.

The paper outline is as follows. In section 2, general expressions for the density of wake charge and total wake charge have been found, using Poisson's equation. In section 3, general expressions for the density of induced surface charges and total charges have been found, using expressions for the normal component of the electric field in the internal and external space of the slab [26]. We apply the results obtained in sections 2 and 3 to the case when slab constructed from a diatomic cubic ionic crystal or polar semiconductor. In section 4, the obtained results are discussed.

## 2 The electromagnetic field of charged particle traversing a slab

We consider a fast particle of charge $q$ moving with a velocity $u$ along the $z$-axis normal to the boundaries of a slab characterized by a local dielectric function $\varepsilon(\omega)$. The time interval $t$ during which the particle moves through the medium is $0 < t < a/u$, where $a$ is the slab thickness. Outside this interval the particle moves in a vacuum.

Ginzburg and Tsytovich [26] have given the expressions for the electromagnetic field of a fast charge passing through a slab. We shall briefly repeat the method of obtaining of these expressions.

Since the problem is homogeneous both in time and directions in each domain $z < 0$, $0 < z < a$, and $z > a$ normal to the charge velocity, it is convenient to represent all field components as Fourier integrals over time and transversal coordinates $\mathbf{r} = (x, y)$. Then the Fourier component of electric field is obtained from the Maxwell equations:

$$\left[\frac{\partial^2}{\partial z^2} + \frac{\omega^2}{c^2}\varepsilon(\omega) - k^2\right]\mathbf{E}(\mathbf{k},\omega,z) = 4\pi\left[-\frac{i\omega u}{c^2}\mathbf{n} + \frac{1}{\varepsilon(\omega)}\left(i\mathbf{k} + \mathbf{n}\frac{\partial}{\partial z}\right)\right]\rho_0(\mathbf{k},\omega,z), \tag{1}$$

where $\mathbf{n} = \mathbf{u}/u$, $\mathbf{k} = (k_x, k_y)$, $\rho_0(\mathbf{k},\omega,z)$ is the Fourier component of the charge density of the particle

$$\rho_0(\mathbf{k},\omega,z) = \frac{q}{(2\pi)^3 u}\exp\left(i\frac{\omega}{u}z\right). \tag{2}$$

The Fourier component of the magnetic field is expressed through $\mathbf{E}(\mathbf{k},\omega,z)$ as follows:

$$\mathbf{B}(\mathbf{k},\omega,z) = \frac{c}{\omega}\left\{-i\left[\nabla \times \mathbf{E}(\mathbf{k},\omega,z)\right] + \left[\mathbf{k} \times \mathbf{E}(\mathbf{k},\omega,z)\right]\right\}. \tag{3}$$

The total solution of (1) for a charge density (2) is a sum of solutions to homogeneous and inhomogeneous equations. While the first equation describes the radiation field, the second equation describes the particle field proper in a medium with local dielectric function $\varepsilon(\omega)$. Also, equation (1) must be solved for each domain inside and outside the slab, and therefore the solutions are joined using the boundary conditions (equality of normal induction components and transverse electric field components on the boundary)

$$E_z(\mathbf{k},\omega,-0) = \varepsilon(\omega)E_z(\mathbf{k},\omega,+0), \quad \varepsilon(\omega)E_z(\mathbf{k},\omega,a-0) = E_z(\mathbf{k},\omega,a+0), \tag{4}$$

$$\mathbf{k}\cdot\mathbf{E}(\mathbf{k},\omega,-0) \quad = \quad \mathbf{k}\cdot\mathbf{E}(\mathbf{k},\omega,+0), \quad \mathbf{k}\cdot\mathbf{E}(\mathbf{k},\omega,a-0) = \mathbf{k}\cdot\mathbf{E}(\mathbf{k},\omega,a+0).$$

Taking account of these conditions, the following system of relations is obtained:

$$\begin{aligned}\mathbf{E}(\mathbf{k},\omega,z) &= \mathbf{E}^{(1)}(\mathbf{k},\omega)\exp\left(i\frac{\omega}{u}z\right)+\frac{2iq}{(2\pi)^2k\omega}a_1^{(-)}\left(k\mathbf{n}+\mathbf{k}\frac{\omega}{kc}\tau_1\right)\exp\left(-i\frac{\omega}{c}\tau_1 z\right), \\ z &< 0\end{aligned} \tag{5}$$

$$\begin{aligned}\mathbf{E}(\mathbf{k},\omega,z) &= \mathbf{E}^{(2)}(\mathbf{k},\omega)\exp\left(i\frac{\omega}{u}z\right)+\frac{2iq}{(2\pi)^2k\omega}\left[a_2^{(-)}\left(k\mathbf{n}+\mathbf{k}\frac{\omega}{kc}\tau_2\right)\exp\left(-i\frac{\omega}{c}\tau_2 z\right)+\right. \\ &\left.+a_2^{(+)}\left(k\mathbf{n}-\mathbf{k}\frac{\omega}{kc}\tau_2\right)\exp\left(i\frac{\omega}{c}\tau_2 z\right)\right], \\ 0 &\le z \le a\end{aligned} \tag{6}$$

$$\begin{aligned}\mathbf{E}(\mathbf{k},\omega,z) &= \mathbf{E}^{(1)}(\mathbf{k},\omega)\exp\left(i\frac{\omega}{u}z\right)+\frac{2iq}{(2\pi)^2k\omega}a_1^{(+)}\left(k\mathbf{n}-\mathbf{k}\frac{\omega}{kc}\tau_1\right)\exp\left(i\frac{\omega}{c}\tau_1 z\right), \\ z &> a\end{aligned} \tag{7}$$

where

$$\mathbf{E}^{(1)}(\mathbf{k},\omega) = -\frac{2iq}{(2\pi)^2}\frac{\omega\mathbf{n}+\gamma^2u\mathbf{k}}{\omega^2+\gamma^2k^2u^2}, \tag{8}$$

$$\mathbf{E}^{(2)}(\mathbf{k},\omega) = -\frac{2iq}{(2\pi)^2}\frac{\omega\left[1-\beta^2\varepsilon(\omega)\right]\mathbf{n}+u\mathbf{k}}{\varepsilon(\omega)\left\{\omega^2\left[1-\beta^2\varepsilon(\omega)\right]+k^2u^2\right\}}, \tag{9}$$

$$\tau_1 = \sqrt{1-k^2c^2/\omega^2}, \quad \tau_2 = \sqrt{\varepsilon(\omega)-k^2c^2/\omega^2}, \tag{10}$$

$$\begin{aligned}a_1^{(-)} &= -\frac{\beta k^2c^2/\omega^2}{\left(1-\beta^2\tau_1^2\right)\left(1-\beta^2\tau_2^2\right)}\frac{1-\varepsilon(\omega)}{D(k,\omega)}\times \\ &\times\left\{f_-^{(1)}\exp\left(i\frac{\omega}{c}\tau_2 a\right)+f_-^{(2)}\exp\left(-i\frac{\omega}{c}\tau_2 a\right)+f_-^{(3)}\exp\left(i\frac{\omega}{u}a\right)\right\},\end{aligned} \tag{11}$$

$$\begin{aligned}a_1^{(+)} &= \frac{\beta k^2c^2/\omega^2}{\left(1-\beta^2\tau_1^2\right)\left(1-\beta^2\tau_2^2\right)}\frac{1-\varepsilon(\omega)}{D(k,\omega)}\exp\left[i\frac{\omega}{u}a\left(1-\beta\tau_1\right)\right]\times \\ &\times\left\{f_+^{(1)}\exp\left(i\frac{\omega}{c}\tau_2 a\right)+f_+^{(2)}\exp\left(-i\frac{\omega}{c}\tau_2 a\right)+f_+^{(3)}\exp\left(-i\frac{\omega}{u}a\right)\right\},\end{aligned} \tag{12}$$

$$f_{\pm}^{(1)} = (\tau_2 - \varepsilon\tau_1)(1 \mp \beta\tau_2)\left(1 \pm \beta\tau_2 - \beta^2\right), \tag{13}$$

$$f_{\pm}^{(2)} = (\tau_2 + \varepsilon\tau_1)(1 \pm \beta\tau_2)\left(1 \mp \beta\tau_2 - \beta^2\right), \tag{14}$$

$$f_{\pm}^{(3)} = 2\tau_2\left[\beta^2\left(1 + \varepsilon - k^2c^2/\omega^2\right) - 1 \pm \beta^2\varepsilon\tau_1\right], \tag{15}$$

$$D(k,\omega) = (\tau_2 + \varepsilon\tau_1)^2 \exp\left(-i\frac{\omega}{c}\tau_2 a\right) - (\tau_2 - \varepsilon\tau_1)^2 \exp\left(i\frac{\omega}{c}\tau_2 a\right) \tag{16}$$

and $\beta = u/c$, $\gamma^{-2} = 1 - \beta^2$. The functions $a_2^{(-)}$ and $a_2^{(+)}$ are expressed through $a_1^{(-)}$ and $a_1^{(+)}$ and are not explicitly given here. They may be obtained from the matching conditions for the normal component of the electric induction on the surfaces $z = 0$ and $z = a$.

The first terms in (6) and (8) describe the Coulomb field of the particle. The first term in (7) describes the particle field in an unbounded medium characterized by the dielectric function $\varepsilon(\omega)$. The field is identical with a Cherenkov radiation electric field in the frequency range $\beta^2\varepsilon(\omega) > 1$. All other terms are due to existence of boundaries. Particularly, they describe the transition radiation in the backward (second term in (6)) and forward (second term in (8)) directions [26].

## 3 The wake charge evaluation

In this section we shall consider the volume charge induced by a moving particle in a slab (the so-called wake charge). To evaluate the wake charge, Poisson equation is used

$$\rho_v = (1/4\pi)\nabla\mathbf{E} - \rho_0 \tag{17}$$

in which the $\rho_0 = q\delta(\mathbf{r})\delta(\xi)$ is the charge density of a test particle, $\xi = z - ut$, $\mathbf{E}$ is the electric field in the slab which is determined by the inverse Fourier transformation of (7). Since the divergence of the second term in (7) (of the radiation field) is zero, the wake-charge density is determined only by the first term in (7).

Using this term in the relation (18) we obtain

$$\rho_v = \frac{q}{2\pi u}\delta(\mathbf{r})\int_{-\infty}^{+\infty} d\omega \exp\left(i\frac{\omega}{u}\xi\right)\frac{1-\varepsilon(\omega)}{\varepsilon(\omega)}, \tag{18}$$

where $\delta(x)$ is a Dirac function. Since the dielectric function of the medium has poles only in the lower $\omega$ half-plane [6], no induced charge exists in front ($\xi > 0$) of the particle. Note that the relation (19) may be also obtained from the expression for electrostatic potential created by the particle in unbounded medium described by a dielectric function $\varepsilon(\omega)$ [6].

Evaluating an integral of expression (19) over the volume we obtain the wake charge, induced by the particle, moving in the slab:

$$Q_v(t) = -q\left[\Phi(t) - \Phi(\tau)\right], \tag{19}$$

where

$$\begin{aligned}\Phi(t) &= \frac{1}{2\pi i}P\int_{-\infty}^{+\infty}\frac{d\omega}{\omega}\exp\left(-i\omega t\right)\frac{1-\varepsilon(\omega)}{\varepsilon(\omega)} = \\ &= -\frac{1}{2}\left(1-\frac{1}{\varepsilon_0}\right)+\theta(t)\left\{1-\frac{1}{\varepsilon_0}-\sum_j \exp\left(-\upsilon_j t\right)\left[A_j\cos\left(\omega_j t\right)+B_j\sin\left(\omega_j t\right)\right]\right\}\end{aligned} \tag{20}$$

and $\tau = t - a/u$, $\varepsilon_0$ is the static dielectric constant of the medium, $\theta(t)$ is the Heaviside unit step function (with $\theta(0) = \frac{1}{2}$), the symbol $P$ denotes the principal value of the integral, $\pm\omega_j - i\upsilon_j$ are the solutions of the equation $\varepsilon(\omega) = 0$ ($\upsilon_j > 0$), while coefficients $A_j$ and $B_j$ are

$$A_j = 2\mathrm{Re}\left\{\frac{1}{\left(\omega_j - i\upsilon_j\right)\varepsilon'\left(\omega_j - i\upsilon_j\right)}\right\}, \quad B_j = 2\mathrm{Im}\left\{\frac{1}{\left(\omega_j - i\upsilon_j\right)\varepsilon'\left(\omega_j - i\upsilon_j\right)}\right\}. \tag{21}$$

Here the prime denotes differentiation with respect to the argument. The summation in (21) is carried over all zeros of the dielectric function.

Analytic properties of the dielectric function [6] and the residue theorem were used in evaluation of $\Phi(t)$.

## 4 Calculation of the induced surface charges

The surface-induced charge density is related to the discontinuity of the electric field $z$-component. Expressions (6)-(8) give

$$\sigma_i(\mathbf{r}, t) = \int d^2\mathbf{k}\int_{-\infty}^{+\infty} d\omega\sigma_i(\mathbf{k}, \omega)\exp[i(\mathbf{kr} - \omega t)], \tag{22}$$

where

$$(2\pi)^2\sigma_0(\mathbf{k},\omega) = \frac{iq}{2\pi\omega}\frac{1-\varepsilon(\omega)}{\varepsilon(\omega)}\left[-\frac{\omega^2}{\omega^2+\gamma^2k^2u^2}+a_1^{(-)}\right], \tag{23}$$

$$(2\pi)^2\sigma_a(\mathbf{k},\omega) = \frac{iq}{2\pi\omega}\frac{\varepsilon(\omega)-1}{\varepsilon(\omega)}\left[-\frac{\omega^2\exp(i\omega a/u)}{\omega^2+\gamma^2k^2u^2}+a_1^{(+)}\exp[i(\omega/c)\tau_1 a]\right] \tag{24}$$

and indices $i = 0$, $a$ refer to the first and second boundaries.

The total induced charge is obtained by integration of expressions (23)-(25):

$$Q_{is}(t) = \int d^2\mathbf{r}\sigma_i(\mathbf{r},t) = (2\pi)^2\int_{-\infty}^{+\infty} d\omega\sigma_i(\omega)\exp(-i\omega t), \tag{25}$$

where $\sigma_i(\omega) = \sigma_i(\mathbf{k}=0,\omega)$.

Let us consider the dynamics of induced charges for the dielectric slab, the dielectric function of which, as it is known, has no singularity in the static limit, when $\omega \to 0$ [6]. It is clear from expressions (12) and (13) that the functions $a_1^{(-)}$ and $a_1^{(+)}$ in expressions (24) and (25) are proportional to $k^2$, and calculating the function $\sigma_i(\omega)$ for the dielectric slab the above-mentioned functions tend to zero. Thus, the charges being induced on the surfaces of the dielectric slab are determined only by the first terms in the expressions (24) and (25), that is by the electric fields which are created by the particle in an unbounded dielectric with the dielectric function $\varepsilon(\omega)$ and in the vacuum when the dielectric is absent.

Using the known Sokhotsky-Plemel relations [6] for quantities $\sigma_i(\omega)$ in (26) we have

$$(2\pi)^2\sigma_0(\omega) = \frac{q}{2\pi i}\frac{1-\varepsilon(\omega)}{\varepsilon(\omega)}P\frac{1}{\omega}, \tag{26}$$

$$(2\pi)^2\sigma_a(\omega) = -\frac{q}{2\pi i}\frac{1-\varepsilon(\omega)}{\varepsilon(\omega)}\exp\left(i\frac{\omega}{u}a\right)P\frac{1}{\omega}. \tag{27}$$

The following relations are obtained from expressions (21), (26)-(28)

$$Q_{os}(t) = q\Phi(t), \quad Q_{as}(t) = -q\Phi(\tau). \tag{28}$$

Thus, expressions (21) and (29) make it possible to obtain the total surface charge as soon as the zeros of the function $\varepsilon(\omega)$ are known. One may easily verify that, at any time, the net induced charge (a sum of $Q_{os}$, $Q_{as}$ and $Q_v$) in the slab is zero.

The following interpretation for expressions (20), (21) and (29) may be given. When the particle approaches the slab surface from a vacuum ($t < 0$) we have $Q_{os} = -(q/2)(1-1/\varepsilon_0)$,

$Q_{as} = (q/2)(1 - 1/\varepsilon_0)$, and $Q_v = 0$. Note that the first boundary of the slab is charged oppositely to the second boundary.

While the particle moves inside the medium ($0 < t < a/u$), the first boundary charge oscillates and decreases:

$$Q_{os}(t) = \frac{q}{2}\left(1 - \frac{1}{\varepsilon_0}\right) - q\sum_j \exp\left(-\upsilon_j t\right)\left[A_j \cos\left(\omega_j t\right) + B_j \sin\left(\omega_j t\right)\right]. \tag{29}$$

Meanwhile the volume charge is increased:

$$Q_v(t) = -q\left(1 - \frac{1}{\varepsilon_0}\right) + q\sum_j \exp\left(-\upsilon_j t\right)\left[A_j \cos\left(\omega_j t\right) + B_j \sin\left(\omega_j t\right)\right]. \tag{30}$$

The second boundary charge remains unchanged.

For $t > a/u$ we have

$$Q_{as}(t) = -\frac{q}{2}\left(1 - \frac{1}{\varepsilon_0}\right) + q\sum_j \exp\left(-\upsilon_j \tau\right)\left[A_j \cos\left(\omega_j \tau\right) + B_j \sin\left(\omega_j \tau\right)\right], \tag{31}$$

$$\begin{aligned} Q_v(t) \quad = \quad & q\sum_j \exp\left(-\upsilon_j t\right)\left[A_j \cos\left(\omega_j t\right) + B_j \sin\left(\omega_j t\right)\right] - \\ & -q\sum_j \exp\left(-\upsilon_j \tau\right)\left[A_j \cos\left(\omega_j \tau\right) + B_j \sin\left(\omega_j \tau\right)\right]. \end{aligned} \tag{32}$$

The charge on the first boundary is given in this case by expression (30).

After the particle crosses the second boundary ($t > a/u$), the charge on the first boundary decreases to its lower limit $(q/2)(1-1/\varepsilon_0)$. On the second boundary the charge value increases and attains its maximum $-(q/2)(1 - 1/\varepsilon_0)$ when $t \gg a/u$. The wake charge in the volume becomes equal to zero.

Thus it follows that after the particle crosses the first boundary, the surface charge is transformed into the wake charge. The latter is transformed again into the surface charge after the particle crosses the second boundary.

We apply the results represented by expressions (20), (21) and (29) for the induced charges in the model of diatomic cubic ionic crystal or polar semiconductor, whose dielectric function is given by [27]

$$\varepsilon(\omega) = \varepsilon_\infty \frac{\omega_L^2 - \omega^2 - i\upsilon\omega}{\omega_T^2 - \omega^2 - i\upsilon\omega}. \tag{33}$$

In this expression $\varepsilon_\infty$ is the optical frequency dielectric function, $\omega_L$ and $\omega_T$ are the frequencies of the longitudinal and transverse-optical vibration modes of infinite wavelength, $\upsilon$ is the damping rate, which we assume to be small ($\upsilon \ll \omega_L$), and $\varepsilon_0$ is the static dielectric function, which enters the theory through the Lyddane-Sachs-Teller relation, $\omega_T^2\varepsilon_0 = \omega_L^2\varepsilon_\infty$. Expression (34) implies that its root $\varepsilon(\omega) = 0$ has the following form:

$$\omega_j - i\upsilon_j = -\frac{i\upsilon}{2} + \Omega, \tag{34}$$

where $\Omega^2 = \omega_L^2 - \upsilon^2/4$. Substituting (34) and (35) into expression (22), we find the coefficients determining the net induced charges:

$$A_j \equiv A = \frac{1}{\varepsilon_\infty} - \frac{1}{\varepsilon_0},\ B_j = gA,\ g = \frac{\upsilon}{2\Omega}. \tag{35}$$

Figure 1 depicts the time dependence of $Q_{os}(t)$, $Q_{as}(t)$, and $Q_v(t)$ for the slab of $LiF$. The following values of the parameters were taken for numerical calculations: $\varepsilon_\infty = 1.96$ and $\varepsilon_0 = 9.01$ [27], $\upsilon/\omega_L = 0.2$, and $\Omega a/u = 15$. We see that as the particle crosses the boundary, the surface and wake charges oscillate with a frequency $\Omega$, although the net induced charge remains equal to zero.

## 5 Conclusion

Let us briefly discuss the conditions in which the processes taking place at the boundaries of the slab can be considered independent and the boundary can be interpreted as that of a half-space, as was done by Gorbunov et al [19].

We see from expressions (20), (21) and (29) that if the condition $a < u/\upsilon_j$ is met, the charge $-q(1-1/\varepsilon_0)$ has no time to transform into the wake charge before the particle reaches the second boundary. For this reason the transformation of a surface charge into a wake charge and the transformation of the latter into a surface charge at the second boundary are interrelated. When the particle crosses the second boundary of the slab, near the boundary it excites electric-field oscillations [3, 20, 26] whose phase is related to that of the oscillations of the electric field near

the first boundary. For the fields as the boundaries to be completely independent, $a$ must exceed $u/v_j$. In this case not only the amplitudes but also the phases of oscillations of the electric fields at the boundaries are independent.

We would like to end our consideration with the following note. The densities of induced surface charges (expressions (23)-(25)) and its corresponding surface current densities may be considered as the sources of transition radiation [27]. As follows from (24) and (25), the densities of induced surface charges are determined by both electrostatic and electromagnetic fields created by the moving particle. The total induced surface charges, determined by (21) and (29), do not depend on the velocity of light and, therefore, are determined only by electrostatic fields. Thus, electromagnetic fields do not contribute into the total induced surface charge. It seems that relativistic effects apparently reveal themselves in the case when the thickness of the medium traversed is limited.

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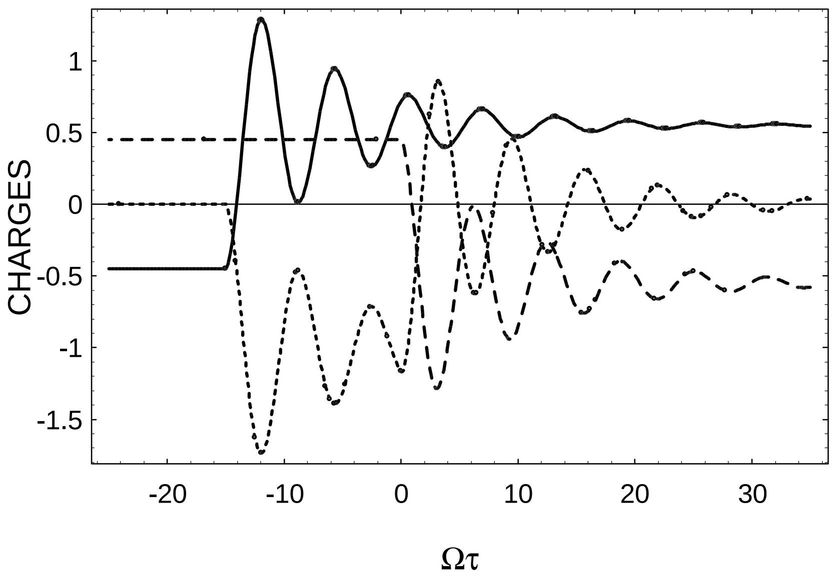
## **Figure Caption**

Fig.1. The dynamics of the induced charge at the front boundary (full curve), the rear boundary (broken curve), and in the volume (dotted curve) for the slab of $LiF$. The following values of the parameters were taken for numerical calculations: $\varepsilon_\infty = 1.96$, $\varepsilon_0 = 9.01$, $v/\omega_L = 0.2$, and $\Omega a/u = 15$.



CHARGES
1
0.5
0
-0.5
-1
-1.5
-20
-10
0
10
20
30
Ωτ